

Space-time description of the hadron interaction at high energies.

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Abstract

In this lecture we consider the strong and electromagnetic interactions of hadrons in a unified way. It is assumed that there exist point-like particles (*partons*) in the sense of quantum field theory and that a hadron with large momentum p consists of $\sim \ln(p/\mu)$ partons which have restricted transverse momenta, and longitudinal momenta which range from p to zero. The density of partons increases with the increase of the coupling constant. Since the probability of their recombination also increases, an equilibrium may be reached. In the lecture we will consider consequences of the hypothesis that the equilibrium really *occurs*. We demonstrate that it leads to constant total cross sections at high energies, and to the Bjorken scaling in the deep inelastic ep-scattering. The asymptotic value of the total cross sections of hadron-hadron scattering turns out to be universal, while the cross sections of quasi-elastic scattering processes at zero angle tend to zero.

The multiplicity of the outgoing hadrons and their distributions in longitudinal momenta (rapidities) are also discussed.

Introduction

In this lecture we will try to describe electromagnetic and strong interactions of hadrons in the same framework which follows from general quantum field theory considerations without the introduction of quarks or other exotic objects.

We will assume that there exist point-like constituents in the sense of quantum field theory which are, however, strongly interacting. It is convenient to refer to these particles as *partons*. We will not be interested in the quantum numbers of these partons, or the symmetry properties of their interactions. We will assume that, contrary to the perturbation theory, the integrals over the transverse momenta of virtual particles converge like in the $\lambda\varphi^3$ theory. It turns out that within this picture a common cause exists for two seemingly very different phenomena: the Bjorken scaling in deep inelastic scattering, and the recent theoretical observation that all hadronic cross sections should approach the same limit (provided

that the Pomeranchuk pole exists). The lecture is organized as follows. In the first part we discuss the propagation of the hadrons in the space as a process of creation and absorption of the virtual particles (partons) and formulate the notion of the parton wave function of the hadron. The second part describes momentum and coordinate parton distributions in hadrons. In the third part we consider the process of deep inelastic scattering. It is shown that from the point of view of our approach the deep inelastic scattering satisfies the Bjorken scaling, and, in contrast to the quark model, the multiplicity of the produced hadrons is of the order of $\ln \frac{\nu}{\sqrt{q^2}}$. The fourth part is devoted to the strong interactions of hadrons and it is shown that in the same framework the total hadron cross sections have to approach asymptotically the same limiting value. In the last part of the lecture we discuss the processes of elastic and quasi-elastic scattering at high energies. It is demonstrated that the cross sections of the quasi-elastic scattering processes at zero angle tend to zero at asymptotically high energies.

Let us discuss, how one can think of the space-time propagation of a physical particle in terms of virtual particles which are involved in the interaction with photon and other hadrons. It is well known that the propagation of a real particle is described by its Green function, which corresponds to a series of Feynman diagrams of the type (for simplicity, we will consider identical scalar particles).

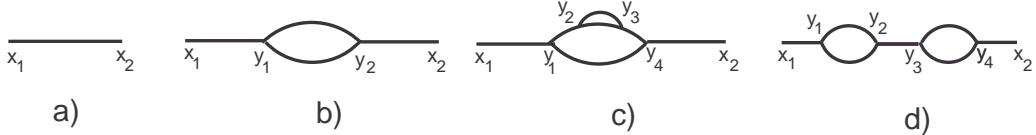


Figure 1:

The Feynman diagrams, having many remarkable properties, have, nevertheless, a disadvantage compared to the old-fashioned perturbation theory. Indeed, they do not show how a system evolves with time in a given coordinate reference frame. For example, depending on the relations between the time coordinates x_{10} , y_{10} , x_{20} and y_{20} , the graph in Fig.1b corresponds to different processes:

Similarly, the diagram Fig.1c corresponds to the processes

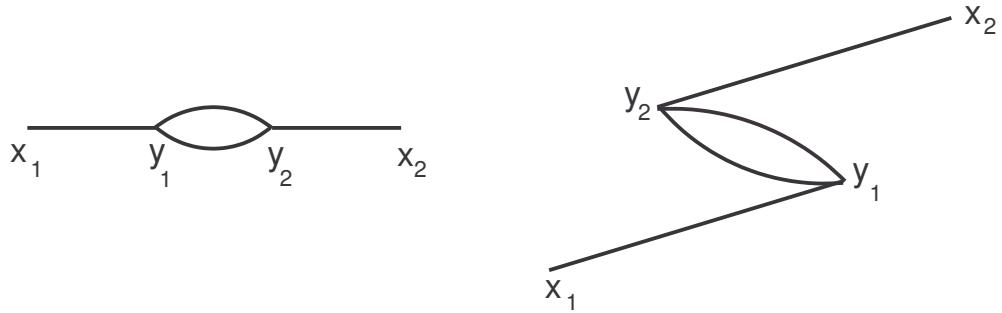


Figure 2:

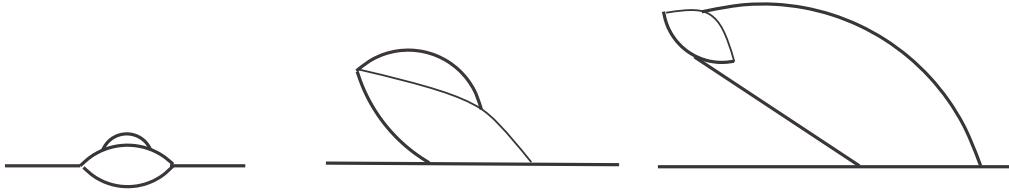


Figure 3:

In quantum electrodynamics, where explicit calculations can be carried out, this complicated correspondence is of little interest. However, for strong interactions, where explicit calculations are impossible, distinguishing between different space developments will be useful.

Obviously, if the interaction is strong (the coupling constant λ is large), many diagrams are relevant. The first question which arises is which configurations dominate: the ones which correspond to the subsequent decays of the particles - the diagrams Fig.2a and Fig3.a, or those which correspond to the interaction of the initial particle with virtual "pairs" created in the vacuum. It is clear that if the coupling constant is large and the momentum of the incoming particle is small (see below), configurations with "pairs" dominate (at least if the theory does not contain infinities). Indeed, if $x_{20} - x_{10}$ is small, then in the case of configurations without "pairs" the integration regions corresponding to each correction will tend to zero with an increase of the number of corrections. At the same time, for the configurations containing "pairs" the region of integration over time will remain infinite. Hence, if the retarded Green function $G^r(y_2 - y_1)$ does not have a strong singularity at $x_{20} - x_{10} \rightarrow 0$, the contribution of the configurations without "pairs" will be relatively small if the coupling constant is large. Even the graphs of the type Fig.1d are determined mainly by configurations like

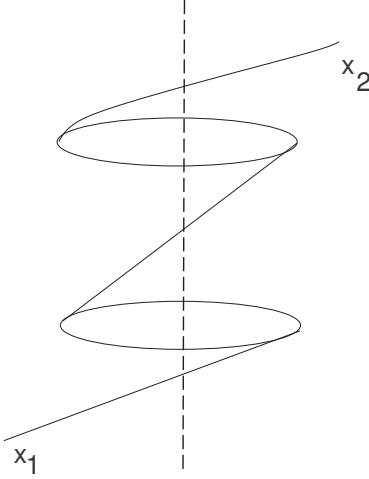


Figure 4:

This means that if we observe a low energy particle at any particular moment of time (the cut in the diagram in Fig. 4), we will see few partons which are decay products of the particle, and a large number of virtual "pairs" which will interact with these partons in the future.

What happens if a particle has a large momentum in our coordinate reference frame? To analyze the space-time evolution of a fast particle we have to consider a space-time interval $(x_1 - x_2)^2$ such that $(x_1 - x_2)^2 \sim \frac{1}{\mu^2}$, and $t_2 - t_1 \sim E/\mu^2 \gg 1/\mu$. Here μ is the mass of the particle, E its energy. In this case, $\vec{x}_1 - \vec{x}_2 = \vec{v}(t_2 - t_1)$, $(x_2 - x_1)^2 = \frac{\mu^2}{E^2}(t_2 - t_1)^2 \sim \frac{1}{\mu^2}$. For such intervals the relation between the configurations with and without "pairs" changes. Configurations corresponding to a decay of one parton into many others start to dominate, while the role of configurations with "pairs" decreases.

The physical origin of this phenomenon is evident. A fast parton can decay, for example, into two fast partons which, due to the energy-time uncertainty relation, will exist for a long time (of the order of E/μ^2), since

$$\Delta E = \sqrt{\mu^2 + \vec{p}^2} - \sqrt{\mu^2 + \vec{p}_1^2} - \sqrt{\mu^2 + (\vec{p} - \vec{p}_1)^2} \sim \frac{\mu^2}{2|\vec{p}|} - \frac{\mu^2}{2|\vec{p}_1|} - \frac{\mu^2}{2|\vec{p} - \vec{p}_1|}.$$

Each of these two partons can again decay into two partons and this will continue up to the point when slow particles, living for a time of the order of $\frac{1}{\mu}$, are created. After that the fluctuations must evolve in the reverse direction, i.e. the recombination of the particles begins.

On the other hand, due to the same uncertainty relation, creation of virtual "pairs" with large momenta in vacuum is possible only for short time intervals of

the order of $\frac{1}{p}$. Hence, it affects only the region of small momentum partons. The way in which this phenomenon manifests itself can be seen using the simplest graph in Fig.5. as an example. We will observe that it is possible to place here many emissions in spite of the fact that the interval x_{12}^2 is of the order of unity ($1/\mu^2$), and the Green function depends only on the invariants.

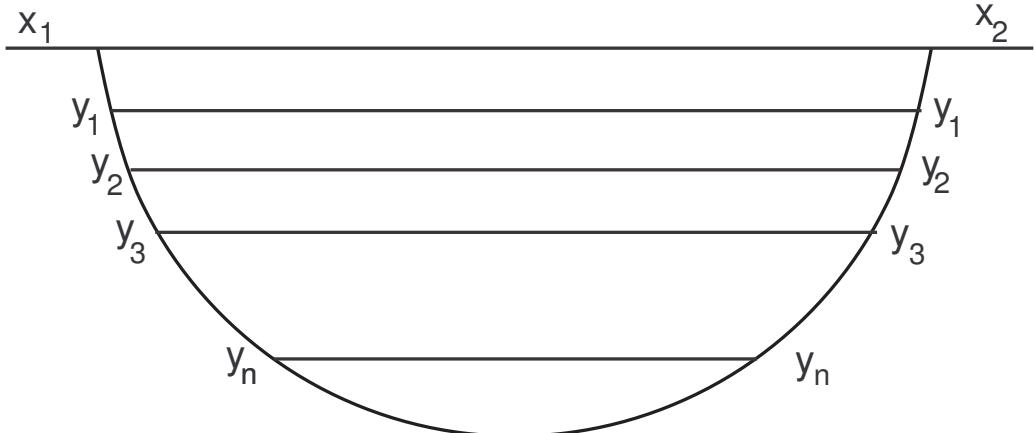


Figure 5:

For the sake of simplicity, let us verify this for one space dimension ($y_i = (t_i, z_i)$). Suppose that $x_1 = (-t, -z)$ and $x_2 = (t, z)$, $x_{12}^2 = (2t)^2 - (2z)^2$. Then $t = z + \frac{x_{12}^2}{8z}$. Let us choose the variables y_i, y'_i in the same way: $y_i = (-t_i, -z_i)$, $y'_i = (t'_i, z'_i)$, and consider the following region of integration in the integral, corresponding to the diagram in Fig.5:

$$1 < z_n < z_{n-1} \dots < z_1 < t,$$

$$1 < z'_n < z'_{n-1} \dots < z'_1 < t,$$

$$z_i \sim z'_i, \quad t_i = z_i + \frac{y_i^2}{2z_i}, \quad t'_i = z'_i + \frac{y'^2_i}{2z'_i}.$$

The integrations over $d^2y_1 \dots d^2y_n d^2y'_1 \dots d^2y'_n$ can be substituted by integrations over y_i^2, y'^2_i and $z_i \equiv y_{iz}, z'_i \equiv y'_{iz}$

$$d^2y_i = \frac{1}{2}dy_i^2 \frac{dz_i}{z_i}, \quad d^2y'_i = \frac{1}{2}dy'^2_i \frac{dz'_i}{z'_i}$$

It is easy to see that in this region of integration the arguments of all Green functions: $(y_i - y'_i)^2, (y_i - y'_{i+1})^2, (y'_i - y'_{i+1})^2$, are of the order of unity, and

the integrals do not contain any small factors. All these conditions for y_i can be satisfied simultaneously for a large number of emissions: $n \sim \ln t$. Indeed, if we write z_n in the form $z_n = C^n$, all conditions will be fulfilled for

$$n \sim \frac{\ln t}{\ln C}, \quad C \geq 1.$$

Obviously, one can consider a more complicated diagram than Fig.5 by including interactions of the virtual particles. On the other hand, configurations containing vacuum "pairs" play a minor role. Moving backwards in time is possible only for short time intervals (Fig.6):

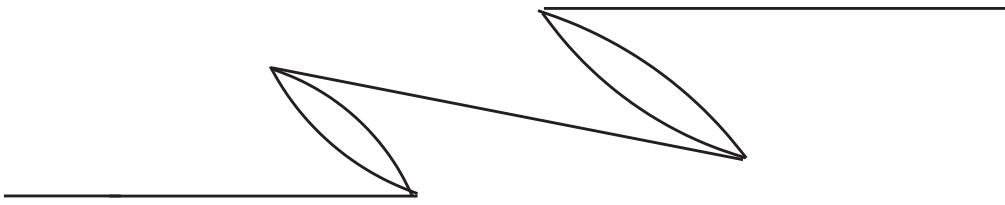


Figure 6:

Hence, we reach the following picture. A real particle with a large momentum p can be described as an ensemble of an indefinite number of partons of the order of $\ln \frac{p}{\mu}$ with momenta in the range from p to zero, and several vacuum pairs with small momenta which in the future can interact with the target.

The observation of a slow particle during an interval of the order of $1/\mu$ does not tell us anything about the structure of the particle since we cannot distinguish it from the background of the vacuum fluctuations, and we can speak only about the interaction of particles or about the spectrum of states. On the contrary, in the case of a fast particle we can speak about its structure, i.e. about the fast partons which do not mix with the vacuum fluctuations. As a result, in a certain sense a fast particle becomes an isolated system which is only weakly coupled to the vacuum fluctuations. Hence, it can be described using a quantum mechanical wave function or an ensemble of wave functions, which determine probabilities of finding certain numbers of partons and their momentum distribution. Such a description is not invariant, since the number of partons depends on the momentum of the particle, but it can be considered as covariant. Moreover, it may be even invariant, if the momentum distribution of the partons is homogeneous in the region of momenta much smaller than the maximal one, and much larger than μ .

Indeed, under the transformation from one reference frame to another in which the particle has, for example, a larger momentum, a new region emerges in the distribution of partons; in the old region, however, the parton distribution remains unchanged. One usually describes hadrons in terms of the quantum mechanics of partons in the reference frame which moves with an infinite momentum, because in this case all partons corresponding to vacuum fluctuations have zero momenta, and such a description is exact. Such a reference frame is convenient for the description of the deep inelastic scattering. However, it is not as good for describing strong interactions, where the slow partons are important. In any case, it appears useful to preserve the freedom in choosing the reference frame and to use the covariant description. This allows a more effective analysis of the accuracy of the derivations.

1 Wave function of the hadron. Orthogonality and normalization

The previous considerations allow us to introduce the hadron wave function in the following way. Let us assume, as usual, that at $t \rightarrow -\infty$ the hadron can be represented as a bare particle (the parton). After a sufficiently long time the parton will decay into other partons and form a stationary state which we call a hadron. Diagrams corresponding to this process are shown in Fig.7.

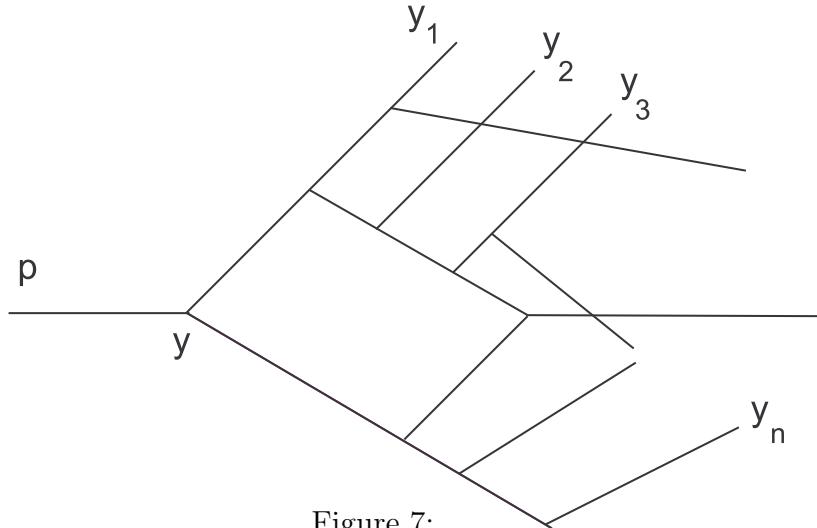


Figure 7:

Let us exclude from the Feynman diagrams those configurations (in the sense of integrations over intermediate times) which correspond to vacuum pair creation.

For the theory $\lambda\varphi^3$ such a separation of vacuum fluctuations corresponds to decomposing φ into positive and negative frequency parts $\varphi = \varphi^+ + \varphi^-$ and sub-

stituting $\varphi^3 = (\varphi^+ + \varphi^-)^3$ by $3(\varphi^{-2}\varphi^+ + \varphi^-\varphi^{+2})$. The previous discussion shows that the ignored term $\varphi^{+3} + \varphi^{-3}$ would mix only partons with small momenta.

It is natural to consider the set of all possible diagrams with a given number of partons n at the given moment of time as a component of the hadron wave function $\Psi_n(t, \vec{y}_1, \dots, \vec{y}_n, p)$. Similarly, we can determine the wave functions of several hadrons with large momenta provided the energy of their relative motion is small compared to their momenta. The latter condition is necessary to ensure that slow partons are not important in the interaction. The Lagrangian of the interaction remains Hermitian even after the terms corresponding to the vacuum fluctuations are omitted. As a result, the wave functions will be orthogonal, and will be normalized in the usual way:

$$\sum_n \int \Psi_n^{b*}(\vec{y}_1, \dots, \vec{y}_n, p_b) i \overleftrightarrow{\partial} \Psi_n^a(\vec{y}_1, \dots, \vec{y}_n, p_a) \frac{d^3 y_1 \dots d^3 y_n}{n!} = (2\pi)^3 \delta(\vec{p}_a - \vec{p}_b) \delta_{ab}, \quad (1)$$

or similarly in the momentum space, after separating

$$\sum_n \frac{1}{n!} \int \Psi_n^{b*}(\vec{k}_1, \dots, \vec{k}_n, \vec{p}) \Psi_n^a(\vec{k}_1, \dots, \vec{k}_n, \vec{p}) \frac{d^3 k_1 \dots d^3 k_n}{2k_{10} \dots 2k_{n0}} \frac{\delta(p - \sum k_i)}{(2\pi)^{3n-1}} = \delta_{ab}. \quad (2)$$

For the momentum range $k_i \gg \mu$, the wave functions coincide with those calculated in the infinite momentum frame. In this reference frame they do not depend on the momentum of the system (except for a trivial factor). This can be easily proven by expanding the parton momenta

$$\vec{k}_i = \beta_i \vec{p} + k_{i\perp}, \quad (3)$$

and writing the parton energy in the form

$$\varepsilon_i = \sqrt{\vec{k}_i^2 + m^2} = \beta_i p + \frac{m^2 + k_{i\perp}^2}{2p\beta_i}. \quad (4)$$

Note now that the integrals which determine Ψ_n , corresponding to Fig.7, can be represented in the form of the old-fashioned perturbation theory where only the differences between the energies of the intermediate states and the initial state $E_k - E$ enter, and the momentum is conserved. Hence, the terms linear in p cancel in these differences, and consequently

$$E_k - E = \frac{1}{2p} \left(\sum_i \frac{m^2 + k_{i\perp}^2}{\beta_i} - m^2 \right) \quad (5)$$

Each consequent intermediate state in Fig.7 in the $\lambda\varphi^3$ model differs from the previous one by the appearance or disappearance of one particle. The factor $\frac{1}{k_i} = \frac{1}{2p} \frac{1}{\beta_i}$, which comes from the propagator of this particle, cancels $2p$ in (5).

Hence, there remain only integrals over $d^2 k_{i\perp} \frac{d\beta_i}{\beta_i}$, and the resulting expression does not depend on p .

$$\sum_n \frac{1}{n!} \int \Psi_n^{b*}(k_{i\perp}, \beta_i) \Psi_n^a(k_{i\perp}, \beta_i) \prod \frac{d^2 k_{i\perp}}{2(2\pi)^2} \frac{d\beta_i}{\beta_i} (2\pi)^3 \delta(1 - \sum \beta_i) = \delta_{ab}. \quad (6)$$

For slow partons, where the expansion (4) is not correct, the dependence on momentum p does not disappear, and contrary to the case of the system moving with $p = \infty$, this dependence cuts off the sum over the number of partons.

2 Distribution of the partons in space and momentum

The distribution of partons in longitudinal momenta can be characterized by the rapidity:

$$\eta_i = \frac{1}{2} \ln \frac{\varepsilon_i + k_{iz}}{\varepsilon_i - k_{iz}}, \quad (7)$$

where k_{iz} is the component of the parton momentum along the hadron momentum.

$$\eta_i \approx \ln \frac{2\beta_i p}{\sqrt{m^2 + k_{i\perp}^2}}. \quad (8)$$

As it is well known, this quantity is convenient since it simply transforms under the Lorentz transformations along the z direction: $\eta'_i = \eta_i + \eta_0$, where η_0 is the rapidity of the coordinate system.

The determination of the parton distribution over η is based on the observation that in each decay process $k_1 \rightarrow k_2 + k_3$ shown in Fig.7 the momenta \vec{k}_2 and \vec{k}_3 are, in the average, of the same order. This means that in the process of subsequent parton emission and absorption the rapidities of the partons change by a factor of the order of unity. At the same time the overall range of parton rapidities is large, of the order of $\ln \frac{2p}{m}$. This implies that in the rapidity space we have short range forces.

Let us consider the density of the distribution in rapidity

$$\varphi(\eta, k_{\perp}, p) = \sum_n \frac{1}{n!} \int |\Psi(k_{\perp}, \eta, k_{\perp 1}, \eta_1, \dots, k_{\perp n}, \eta_n)|^2 (2\pi)^3 \delta(\vec{p} - \vec{k} - \sum \vec{k}_i) \prod \frac{dk_i d\eta_i}{2(2\pi)^3} \quad (9)$$

in the interval $1 \ll \eta \ll \eta_p$ (see Fig.8).

The independence of φ on p for these values of η means that Ψ depends only on the differences $\eta_i - \eta_p$. If $\varphi = \varphi(\eta - \eta_p, k_{\perp})$ decreases with the increase of $\eta - \eta_p$, this corresponds to a weak coupling, i.e. to a small probability of the decay of the initial parton. If the coupling constant grows, the number of partons increases

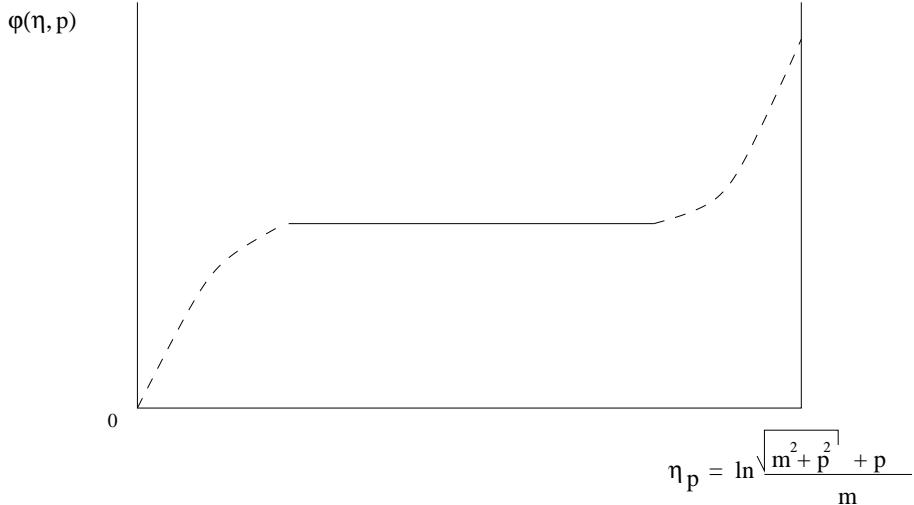


Figure 8:

and at a *certain* value of the coupling constant an equilibrium is reached, since the probability of recombination also increases. The value of this critical coupling constant has to be such that the recombination probability due to the interaction should be larger than the recombination probability related to the uncertainty principle.

The basic hypothesis is that such an equilibrium *does occur* and that due to the short-range character of interaction it is *local*. This is equivalent to the hypothesis of the constant total cross sections of interaction at $p \rightarrow \infty$. Hence we assume that the equilibrium is determined by the vicinity of the point η of the order of unity and it does not depend on η_p . Obviously, this can be satisfied only if $\varphi(\eta, \eta_p, k_\perp) = \varphi(k_\perp)$ does not depend on η and η_p at $1 \ll \eta \ll \eta_p$. According to the idea of Feynman, this situation resembles the case of a sufficiently long one-dimensional matter in which, due to the homogeneity of the space, far from the boundaries the density is either constant or oscillating (for a crystal). In our case the analogue of the homogeneity of space is the relativistic invariance (the shift in the space of rapidities). *For the time being* we will not consider the case of the crystal. According to (9), the integral of $\varphi(\eta, \eta_p, k_\perp)$ over η and k_\perp has the meaning of the average parton density which is, obviously, of the order of $\eta_p \sim \ln \frac{2p}{m}$.

Generally speaking, we cannot say anything about the parton distribution in the transverse momenta except for one statement: it is absolutely crucial for the whole concept that it must be restricted to the region of the order of parton masses, like in the $\lambda\varphi^3$ theory.

Consider now the spatial distribution of the partons. First, let us discuss

parton distribution in the plane perpendicular to the momentum \vec{p} . For that purpose it is convenient to transform from $\Psi_n(k_{1\perp}, \eta_1, k_{2\perp}, \eta_2, \dots, k_{n\perp}, \eta_n)$ to the impact parameter representation $\Psi_n(\vec{\rho}_1, \eta_1, \vec{\rho}_2, \eta_2, \dots, \vec{\rho}_n, \eta)$:

$$\Psi_n(\vec{\rho}_n, \eta_n) = \int e^{i \sum k_{i\perp} \rho_i} \Psi(k_{i\perp}, \eta_i) \delta(\sum k_{\perp i}) (2\pi)^2 \prod \frac{d^2 k_i}{(2\pi)^2}. \quad (10)$$

Let us rank the partons in the order of their decreasing rapidities. Consider a parton with the rapidity $\eta \ll \eta_p$ and let us follow its history from the initial parton. Initially, we will assume that it was produced solely via parton emissions (Fig.9).

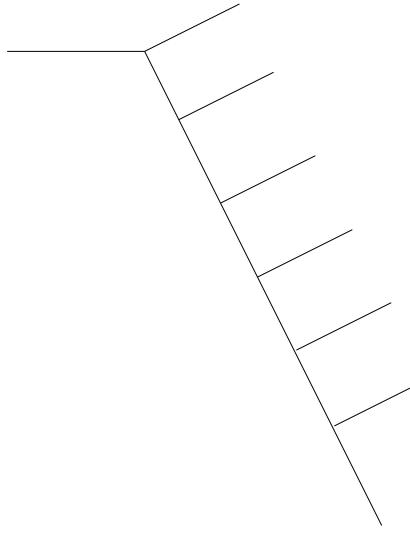


Figure 9.

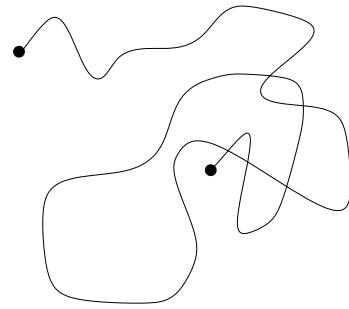


Figure 10.

In this case it is clear that if the transversal momenta of all partons are of the order of μ , than each parton emission leads to a change of the impact parameter $\vec{\rho}$ by $\sim \frac{1}{\mu}$. If n emissions are necessary to reduce the rapidity from η_p to η , and they are independent and random, $\overline{(\Delta \rho)^2} \sim n$. If every emission changes the rapidity of the parton by about one unit, then

$$\overline{(\Delta \rho)^2} = \gamma(\eta_p - \eta). \quad (11)$$

Hence, the process of the subsequent parton emissions results in a kind of diffusion in the impact parameter plane. The parton distribution in ρ for the rapidity η

has the Gaussian form:

$$\varphi(\rho, \eta) = \frac{C(\eta)}{\pi\gamma(\eta_p - \eta)} e^{-\frac{\rho^2}{\gamma(\eta_p - \eta)}}, \quad (12)$$

if the impact parameter of the initial parton is considered as the origin. Consequently, the partons with $\eta \approx 0$ have the broadest distribution, and, hence, the fast hadron is of the size

$$R = \sqrt{\gamma\eta_p} \approx \sqrt{\gamma \ln \frac{2p}{m}}. \quad (13)$$

The account of the recombination and the scattering of the partons affects only densities of partons and fluctuations, but does not change the radius of the distribution which can be viewed as the front of the diffusion wave.

Let us discuss the parton distribution over the longitudinal coordinate. A relativistic particle with a momentum p is commonly considered as a disk of the thickness $1/p$. In fact, this is true only in the first approximation of the perturbation theory. Really, a hadron is a disk with radius $\sqrt{\gamma \ln \frac{2p}{m}}$ and the thickness of the order of $1/\mu$. Indeed, each parton with a longitudinal momentum k_{iz} is distributed in the longitudinal direction in an interval $\Delta z_i \sim \frac{1}{k_{iz}}$. Since the parton spectrum exists in the range of momenta from p down to $k_i \sim \mu$, the longitudinal projection of the hadron wave function has the structure depicted in Fig.11.

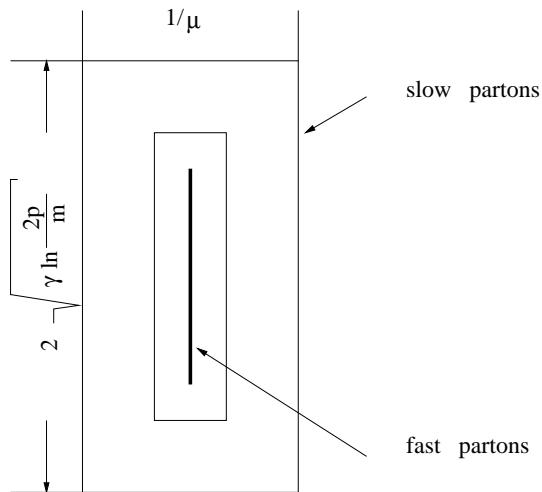


Figure 11:

Finally, let us consider what is the lifetime of a particular parton. As we have discussed in the Introduction, in a theory which is not singular at short distances,

the intervals y_{12}^2 between two events represented by a Feynman diagram are of the order of unity. For a fast particle moving along the z axis, $z_{21} = vt_{21}$ and $y_{12}^2 = t_{21}^2 \frac{m^2}{p^2}$. Consequently, the lifetime of a fast parton with a momentum k_i is of the order of $\frac{k_i}{\mu^2}$. The presented arguments were based on the $\lambda\varphi^3$ theory which is the only theory which provides a cutoff in transverse momenta. Still, the argument should hold for other theories and for particles with spins, if one assumes that in these theories the cutoff of transverse momenta occurs in some way. On the other hand, the $\lambda\varphi^3$ theory cannot be considered as a self-consistent example. Indeed, due to the absence of a vacuum state, the series of perturbation theory do not make sense (series with positive coefficients are increasing as factorials). Hence, the picture we have presented here does not correspond *literally* to any particular field theory. At the same time, it corresponds fully to the main ideas of the quantum field theory and to its basic space-time relations.

3 Deep inelastic scattering

It is convenient to consider the deep inelastic scattering of electrons in the frame where the time component of the virtual photon momentum is $q_0 = 0$. In this reference frame the momentum of the photon is equal to $-q_z$ ($q^2 = -q_z^2$), while the momentum of the hadron is $p_z = \omega q_z/2$ ($\omega = -2p \cdot q/q^2$). Suppose that q_z is large and $\omega \sim 1$. According to our previous considerations, a fast hadron can be viewed as an ensemble of partons. In this system a photon looks as a static field with the wavelength $\sim 1/q_z$.

The main question is, with which partons can the photon interact. We can consider the static field of a photon as a packet with a longitudinal size of the order of $1/q_z$. The interaction time between a hadron with the size $1/\mu$ and such a packet is of the order of $1/\mu$. However, due to the big difference between the parton and photon wave lengths, the interaction with a slow parton is small. Hence, the photon interacts with partons which have momenta of the order of q_z . Partons with such momenta are distributed in the longitudinal direction in the region $1/q_z$. Because of this, the time of the hadron-photon interaction is in fact of the order of $1/q_z$, i.e. much shorter than the lifetime of a parton. This means that the photon interacts with a parton as with a free particle, and so not only the momentum but also the energy is conserved. As a result, the energy-momentum conservation laws select the parton with momentum $q_z/2$, which can absorb a photon

$$k_{iz} - q_z = k'_{iz}, \quad |k_{iz} - q_z| = k_{iz}.$$

This gives

$$k_{iz} = \frac{q_z}{2}, \quad k'_{iz} = -\frac{q_z}{2}.$$

The cross section of such a process is, obviously, equal to the cross section σ_0 of the absorption of a photon by a free particle, multiplied by the probability to

find a parton with a longitudinal momentum $q_z/2$ inside the hadron, i.e. by the value $\varphi(\eta_{q/2}, \eta_p)$ (9), integrated over k_\perp . (The necessary accuracy of fulfilment of the conservation laws allows any $k_\perp \ll q_z$).

As it was already discussed, $\varphi(\eta, \eta_p) = \varphi(\eta - \eta_p) \equiv \varphi(\omega)$. Hence, using the known cross section for the interaction of the photon with a charged spinless particle, we obtain for the cross section of the deep inelastic scattering

$$\frac{d^2\sigma}{dq^2 d\omega} = \frac{4\pi\alpha^2}{q^4} \left(1 - \frac{pq}{pp_e} \right) \varphi(\omega), \quad (14)$$

where p_e is the electron momentum. If the partons have spins, the situation becomes more complicated, since the cross sections of the interactions between photons and partons with different spins are different. The parton distributions in rapidities for different spins may also be different, leading to the form:

$$\frac{d^2\sigma}{dq^2 d\omega} = \frac{4\pi\alpha^2}{q^4} \left\{ \left(1 - \frac{(pq)}{(pp_e)} \right) \varphi_0(\omega) + \left[1 - \frac{pq}{pp_e} + \frac{1}{2} \left(\frac{pq}{pp_e} \right)^2 \right] \varphi_{\frac{1}{2}}(\omega) \right\}. \quad (15)$$

Let us discuss now a very important question, namely: what physical processes take place in deep inelastic scatterings. To clarify this, we go back to Fig.7 determining the hadron wave function. We will neglect the parton recombinations in the process of their creation from the initial parton, i.e. we consider fluctuations of the type shown in Fig.9. Suppose that the photon was absorbed by a parton with a large momentum $q_z/2$. As a result, this parton obtained a momentum $-q_z$ and moves in the opposite direction with momentum $-q_z/2$. The process is depicted in Fig.12. What will now happen to this parton and to the remaining partons? Within the framework we are using it is highly unlikely that the parton with the momentum $-q_z/2$ will have time to interact with the other partons. The probability to interact directly with residual partons will be small, because the relative momentum of the parton with $-q_z/2$ and the rest of the partons is large. It could interact with other partons after many subsequent decays which, in the end, could create a slow parton. However, the time needed for these decays is large, and during this time the parton and its decay products will move far away from the remaining partons, thus the interaction will not take place.

Hence, we come to the conclusion that one free parton is moving in the direction $-q_z$. What will we observe experimentally, if we investigate particles moving in this direction? To answer this question, it is sufficient to note that, in average, a hadron with a momentum k_z consists of n partons; $n = C \ln \frac{k_z}{\mu}$ at $k_z \gg \mu$.

In a sense there should exist an uncertainty relation between the number of partons in a hadron (n) and the number of hadrons in a parton (n_p)

$$n_p n \gtrsim c \ln \frac{k_z}{\mu}, \quad (16)$$

where k_z is the momentum of the state.

We came to the conclusion that the parton decays into a large number of hadrons i.e. in fact the parton is very short-lived, highly virtual. Hence, we have to discuss whether this conclusion is consistent with the assumption that the photon-parton interaction satisfies the energy conservation. To answer this question, let us calculate the mass of a virtual parton with momentum k_z , decaying into n hadrons with momenta k_i and masses m_i .

$$M^2 = (\sum \sqrt{m_i^2 + k_i^2})^2 - k_z^2 = \left(k_z + \sum_i \frac{m_i^2 + k_{i\perp}^2}{2k_{iz}} \right)^2 - k_z^2 \approx k_z \sum_i \frac{m_i^2 + k_{i\perp}^2}{k_{iz}}.$$

If the hadrons are distributed almost homogeneously in rapidities, their longitudinal momenta decrease exponentially with their number, and in the sum only a few terms, corresponding to slow hadrons, are relevant. As a result, $M^2 \sim k_z \mu$, i.e. the time of the existence of the parton is of the order of $1/\mu$, much larger than the time of interaction with a photon $1/q_z$. Let us discuss now, what happens to

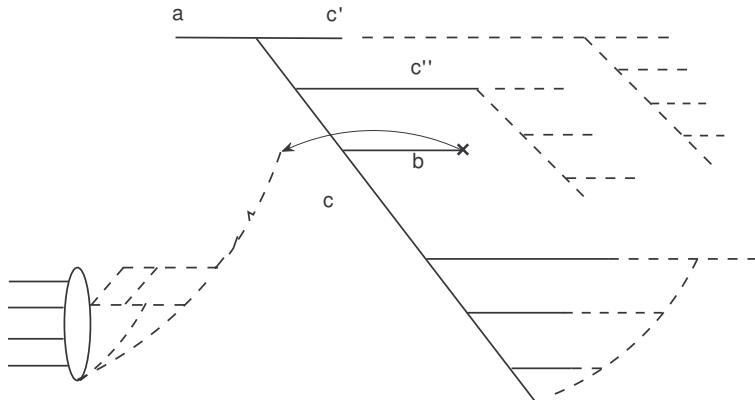


Figure 12:

the remaining partons. Little can be determined using only the uncertainty relation eq.(16). This is because the number of partons before the photon absorption was n , after the photon absorption it became $n - 1$ and, consequently, according to the uncertainty relation, the number of hadrons corresponding to this state can range from 1 to n . Hence, everything depends on the real perturbation of the hadron wave function due to the photon absorption.

Consider now the fluctuation shown in Fig.12. The photon absorption will not have any influence on partons created after the parton "b" which absorbed the photon was produced, and which have momenta smaller than "b". These

fluctuations will continue, and the partons can, in particular, recombine back into the parton "c". The situation is different for partons which occurred earlier and have large momenta ("c'", "c'"). In this case the fluctuation cannot evolve further the same way, since the parton "b" has moved in the opposite direction. As a result, it is highly probable that partons "c'" and "c'" will move apart and lose coherence. On the other hand, slow partons which were emitted by "c'" and "c'" earlier and which are not connected with the parton "b", will be correlated, as before, with each of them. Thus "c'" and "c'" will move in space together with their slow partons, i.e. in the form of hadrons. Hence, it appears that partons flying in the initial direction lead to the production of the order of $c \ln \frac{\omega q_z}{2} - c \ln \frac{q_z}{2} = c \ln \omega$ hadrons with rapidities ranging from $\ln \frac{\omega q_z}{\mu}$ to $\ln \frac{q_z}{\mu}$. This answer can be interpreted in the following way. After the photon is absorbed, a hole is created in the distribution of partons moving in the initial direction.

Contrary to the case of rapidities of partons, we will count the rapidity of the hole not from zero rapidity but from the rapidity $\ln \frac{\omega q_z}{\mu}$. In this case the rapidity of the hole is $\ln \omega$. If we now represent the parton hole with rapidity $\ln \omega$ as a superposition of the hadron states, this superposition will contain $\ln \omega$ hadron states.

Let us represent the whole process by a diagram describing rapidity distributions of partons and hadrons. Before the photon absorption the partons in the hadrons are distributed at rapidities between zero and $\ln \frac{\omega q_z}{\mu}$, while after the photon absorption a parton distribution is produced which is shown in Fig. 13.

This parton distribution leads to the hadron distribution shown in Fig. 14. The total multiplicity corresponding to this distribution is

$$\bar{n} = c \ln \frac{q_z}{\mu} + c \ln \omega = c \ln \frac{\nu}{\mu \sqrt{-q^2}}.$$

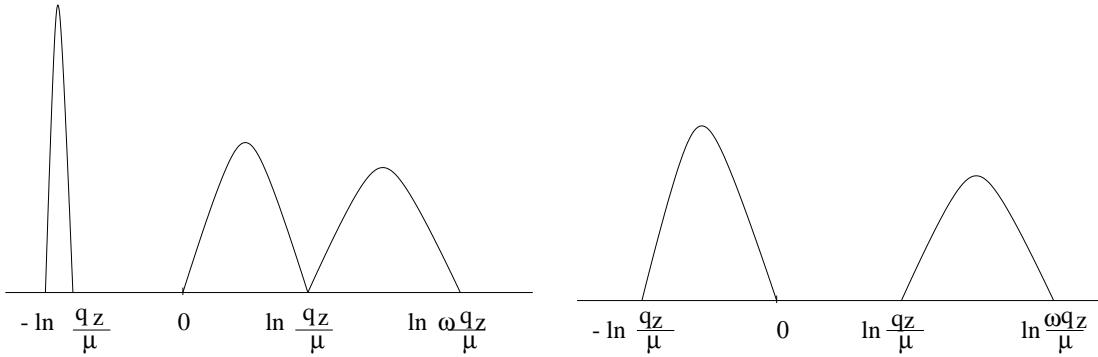


Figure 13

Figure 14

This hadron distribution in rapidities in the deep inelastic scattering differs

qualitatively from those previously discussed in the literature. It corresponds to $c \ln \frac{\sqrt{-q^2}}{\mu}$ hadrons moving in the photon momentum direction, and $\ln \omega$ hadrons are moving in the nucleon momentum direction, with a gap in rapidity between these distributions. The hadron distribution which was obtained in the framework of perturbation theory for superconverging theories like $\lambda \varphi^3$ (Drell, Yan) differs qualitatively from the distribution in Fig. 14.

In conclusion of this part, it is necessary to point out that the problem of spin properties of the partons exist in this picture even if the partons do not have quark quantum numbers. If, as the experiment shows, the cross section σ_T for the interaction of the transversal photons is larger than the cross section for the interaction of the longitudinal photons, σ_L , the charged partons have predominantly spin 1/2. This means that at least one fermion, for example a nucleon, has to move in the direction of the photon momentum. In other words, in deep inelastic scattering the distribution of the created hadrons in quantum numbers as the function of their rapidities differs essentially from what we are used to in the strong interactions. Perhaps this is one of the key prediction of the non-quark parton picture for $\sigma_T \gg \sigma_L$.

4 Strong interactions of hadrons

Let us discuss now the strong interactions of hadrons. First, we consider a collision of two hadrons in the laboratory frame. Suppose that a hadron "1" with momentum \vec{p}_1 hits hadron "2" which is at rest. Obviously, the parton wave function makes no sense for the hadron at rest, since for the latter the vacuum fluctuations are absolutely essential. However, the hadron at rest can also be understood as an ensemble of slow partons distributed in a volume of the order of $1/\mu$, independent of the origin of the partons. Indeed, it does not matter whether these partons are decay products of the initial parton or the result of the vacuum fluctuations. How can a fast hadron, consisting of partons with rapidities from $\ln \frac{2p_1}{\mu}$ to zero, interact with the target which consists of slow partons? Obviously, the cross section of the interaction of two point-like particles with a large relative energy is not larger than $\pi \lambda^2 \sim \frac{1}{s_{12}} \sim e^{-\eta_{12}}$ (where λ is the wave length in the c.m. frame, η_{12} is the relative rapidity). That is why only slow partons of the incident hadron can interact with the target with a cross section which is not too small. This process is shown in Fig. 15.

If the slow parton which initiated the interaction was absorbed in this interaction, the fluctuation which lead to its creation from a fast parton was interrupted. Hence, all partons which were emitted by the fast parton in the process of fluctuation cannot recombine any more. They disperse in space and ultimately decay into hadrons leading to the creation of hadrons with rapidities from zero to $\ln \frac{2p_1}{\mu}$. The interaction between the partons is short-range in rapidities. Hence, the hadron distribution in rapidities will reproduce the parton distribution in

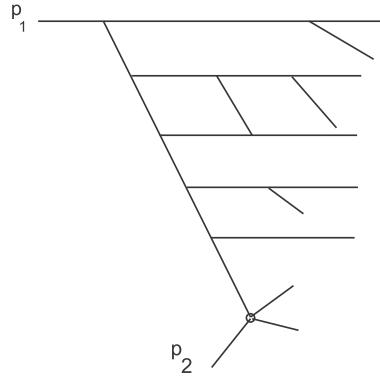


Figure 15:

rapidities. In particular, the inclusive spectrum of hadrons will have the form shown in Fig. 8, with an unknown distribution near the boundaries. The total hadron multiplicity will be of the order of $\eta_p = \ln \frac{2p_1}{\mu}$. If the probability of finding a slow parton in the hadron does not depend on the hadron momentum (this would be quite natural, since with the increase of the momentum the life-time of the fluctuation is also growing), the total cross section of the interaction will not depend on the energy at high energies.

Before continuing the analysis of inelastic processes, let us discuss, how to reconcile the energy independence of the total interaction cross section at high energies with the observation discussed above that the transverse hadron sizes increase with the increase of the energy as $\sqrt{\gamma \ln \frac{2p}{\mu}}$. The answer is that slow partons are distributed almost homogeneously over the disk of the radius $\sqrt{\gamma \ln \frac{2p}{\mu}}$ (Eq.(11)) , while their overall multiplicity during the time of $1/\mu$ is of order of unity.

Let us see now how the same process will look, for example, in the c.m. frame. In this reference frame the interaction will have the form as shown in Fig. 16.

Each of the hadrons consists of partons with rapidities ranging from $-\ln \frac{2p_c}{\mu}$ to zero and from zero to $\ln \frac{2p_c}{\mu}$, respectively. The slow partons interact with cross sections which are not small. As a result, the fluctuations will be interrupted in both hadrons, and the partons will fly away in the opposite directions, leading to the creation of hadrons with rapidities from $-\ln \frac{2p_c}{\mu}$ to $\ln \frac{2p_c}{\mu}$. From the point of view of this reference frame the inclusive spectrum must have the form shown in Fig. 17, with unknown distributions not only at the boundaries but also in the centre, since the distribution of the slow partons in the hadrons and in vacuum fluctuations is unknown. The hadron inclusive spectrum, however, should not depend on the reference frame. Thus the inclusive spectrum in Fig. 17 should

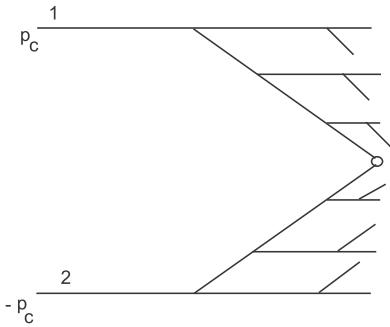


Figure 16.

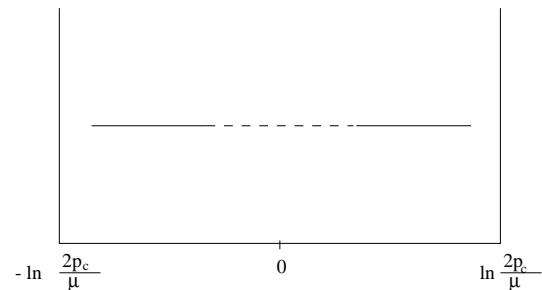


Figure 17

coincide with the inclusive spectrum in Fig. 8, and they should differ only by a trivial shift along the rapidity axis, i.e. due to relativistic invariance we know something about the spectra of slow partons and vacuum fluctuations.

Let us demonstrate that this comparison of processes in two reference frames leads to a very important statement, namely that at ultra-high energies the total cross sections for the interactions of arbitrary hadrons should be equal. Indeed, we have assumed that the distribution of hadrons reproduces the parton distribution.

From the point of view of the laboratory frame the distribution of partons and, consequently, distribution of hadrons in the central region of the spectrum is completely determined by the properties (quantum numbers, mass, etc.) of particle 1, and does not depend on the properties of particle 2. On the other hand, from the point of view of the antilaboratory frame (where the particle 1 is at rest) everything is determined by the properties of particle 2. This is possible only if the distribution of partons in the hadrons with rapidities η much smaller than the hadron rapidity η_p does not depend on the quantum numbers and the mass of the hadron, that is the parton distribution with $\eta \ll \eta_p$ should be universal. From the point of view of the c.m. system the same region is determined by slow partons of both hadrons and by vacuum fluctuations (which are universal), and, consequently, the distribution of slow partons is also universal.

It is natural to assume that the probability of finding a hadron in a sterile state without slow partons tends to zero with the increase of its momentum, in other words assume that slow partons are always present in a hadron (compare to the decrease of the cross section of the elastic electron scattering at large q^2). In this case considering the process in the c.m. system, we see that the total cross section of the hadron interaction is determined by the cross section of the interaction of slow partons and by their transverse distribution which is universal. Consequently, the total hadron interaction cross section is also universal, i.e. equal for any hadrons.

This statement looks rather strange if we regard it, for instance, from the fol-

lowing point of view. Let us consider the scattering of a complicated system with a large radius, for example, deuteron-nucleon scattering. As we know, the cross section of the deuteron-nucleon interaction equals the sum of the nucleon-nucleon cross sections, thus it is twice as large as the nucleon-nucleon cross section. How and at what energies can the deuteron-nucleon cross section become equal to the nucleon-nucleon cross section? How is it possible that the density of slow partons in the deuteron turns out to be equal to the density of slow partons in the nucleon? To answer this question, let us discuss the parton structure of two hadrons which are separated in the plane transverse to their longitudinal momenta by a distance much larger than their Compton wave length $1/\mu$. Suppose that at the initial moment they were point-like particles. Next, independently of each other, they begin to emit partons with decreasing longitudinal momenta. At the same time the diffusion takes place in the transverse plane so that the partons will be distributed in a growing region. The basic observation which we shall prove and which answers our question is that if the momenta of the initial partons are sufficiently large, then during one fluctuation the partons coming from different initial partons will inevitably meet in space (Fig. 18) in the region of the order of $1/\mu$. They will have similar large rapidities and, hence, will be able to interact with a probability of the order of unity. If such “meetings” take place sufficiently frequently, the probability of the parton interaction will be unity. Consequently, the further evolution and the density of the slow partons which are created after the meeting may not depend on the fact that initially the transverse distance between two partons was large.

In terms of the diffusion in the impact parameter plane this statement corresponds to the following picture. Suppose that initial partons were placed at points ρ_1 and ρ_2 in Fig. 19 and that their longitudinal momenta are of the same order of magnitude, i.e. difference of their rapidities is of the order of unity, while each of the rapidities is large. We will follow the parton starting from point ρ_1 , which decelerates via emission of other partons. As we have seen, its propagation in the perpendicular plane corresponds to diffusion. The difference of rapidities $\eta_p - \eta$ at the initial and the considered moments serves the role of time in this diffusion process.

The diffusion character of the process means that the probability density of finding a parton with rapidity η at the point ρ if it started from the point ρ_1 with rapidity η_p is

$$\omega(\vec{\rho}, \vec{\rho}_1, \eta_p - \eta) = \frac{1}{\pi \gamma(\eta_p - \eta)} e^{-\frac{(\vec{\rho} - \vec{\rho}_1)^2}{\gamma(\eta_p - \eta)}}. \quad (17)$$

The situation is exactly the same for a decelerating parton which started from the point ρ_2 . Thus, the probability of finding both partons at the same point ρ with equal rapidities is proportional to

$$\omega(\rho_{12}, \eta_p - \eta) =$$

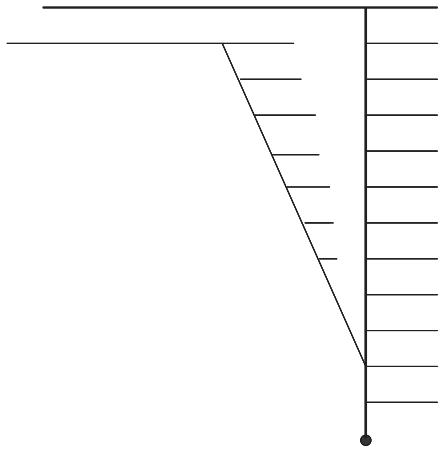


Figure 18

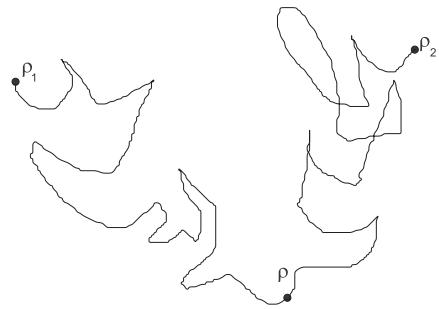


Figure 19

$$\int \omega(\vec{\rho}, \vec{\rho}_1, \eta_p - \eta) \omega(\vec{\rho}, \vec{\rho}_2, \eta_p - \eta) d^2 \rho = \\ \frac{1}{2\pi\gamma(\eta_p - \eta)} \exp \left[-\frac{(\vec{\rho}_1 - \vec{\rho}_2)^2}{2\gamma(\eta_p - \eta)} \right] \quad (18)$$

If we now integrate this expression over η , i.e. estimate the probability for the partons to meet at some rapidities, we obtain

$$\int_0^{\eta_p} \omega(\rho_{12}, \eta_p - \eta) d\eta \approx \frac{1}{\pi} \log \frac{2\gamma\eta_p}{\rho_{12}^2} \Big|_{\eta_p \rightarrow \infty} \longrightarrow \infty. \quad (19)$$

This means that if $2\gamma\eta_p \gg \rho_{12}^2$, the partons will inevitably meet. According to (19) we get a probability much larger than unity. The reason is that under these conditions the meeting of partons at different values of η are not independent events and therefore it does not make sense to add the probabilities. It is easy to prove this statement directly, for example with the help of the diffusion equation. We will not do this, however. According to a nice analogy suggested by A. Larkin, this theorem is equivalent to the statement that if you are in an infinite forest in which there is a house on a finite distance from you, then, randomly wandering in the forest, you sooner or later arrive at this house. Essentially, the reason is that in the two-dimensional space the region inside of which the diffusion takes place and the length of the path travelled during the diffusion increase with time in the same way. From the point of view of the reference frame in which the

deuteron is at rest and is hit by a nucleon in the form of a disk, the radius of which is much larger than that of the deuteron, the statement of the equality of cross sections means that the parton states inside the disk are highly coherent.

It is clear from above that the cross sections of two hadrons can become equal only when the radius of parton distribution $\sqrt{\gamma \eta_p}$ which is increasing with the energy becomes much larger than the size of both hadrons. Substituting $4 \cdot \frac{0.3}{m^2}$ for the value of γ (m is the proton mass)¹ we see that the deuteron-nucleon cross section will practically never coincide with the nucleon-nucleon cross section, while the tendency for convergence of cross sections for pion-nucleon, kaon-nucleon and nucleon-nucleon scatterings may be manifested already starting at the incident energies $\sim 10^3$ GeV.

5 Elastic and quasi-elastic processes

So far we focused on the implications of the considered picture for inelastic processes with multiplicities, growing logarithmically with the energy. However, with a certain probability it can happen that slow partons scatter at very small angles and the fluctuations will not be interrupted in either of the hadrons (for example, if we discuss the process in the c.m. frame). In this case small angle elastic or quasi-elastic scattering will take place (Fig. 20). First, let us calculate the elas-

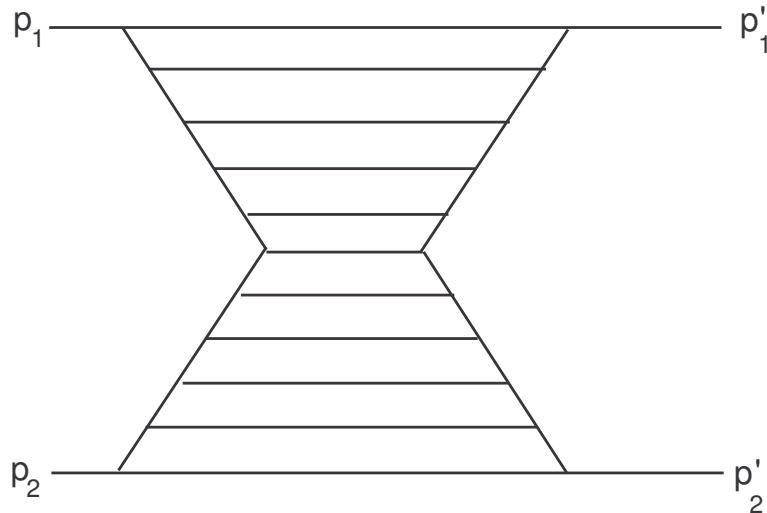


Figure 20:

tic scattering amplitude. It is well known that the imaginary part of the elastic

¹It will be demonstrated below that $\gamma = 4\alpha'$, where α' is the slope of the Pomeron trajectory. The current data give $\alpha' \sim 0.25 \text{GeV}^{-2}$.

scattering amplitude can be written in the form

$$A_1(s_{12}) = s_{12} \int d^2 \rho_{12} e^{i\vec{q}\vec{\rho}_{12}^2} \sigma(\rho_{12}, s_{12}), \quad (20)$$

where s_{12} is the energy squared in the c.m. system, ρ_{12} is the relative impact parameter, $\sigma(\rho_{12}, s_{12})d^2\rho_{12}$ - the total interaction cross section of particles being at the distance ρ_{12} and \vec{q} is the momentum transferred. In order to calculate $\sigma(\rho_{12}, s_{12})$ it is sufficient to notice that, according to (12), the probability of finding a slow parton with rapidity η at the impact parameter ρ'_1 which originated from the first hadron with an impact parameter $\vec{\rho}_1$ is

$$\varphi_1(\vec{\rho}_1, \vec{\rho}'_1, \eta_1, \eta_{pc}) \frac{C(\eta_1)}{\pi \gamma \eta_{pc}} \exp \left[-\frac{(\vec{\rho}_1 - \vec{\rho}'_1)^2}{\gamma \eta_{pc}} \right]. \quad (21)$$

The probability a parton originating from the second hadron at impact parameter ρ'_2 is

$$\varphi_2(\vec{\rho}_2, \vec{\rho}'_2, \eta_2, \eta_{pc}) \frac{C(\eta_2)}{\pi \gamma \eta_{pc}} \exp \left[-\frac{(\vec{\rho}_2 - \vec{\rho}'_2)^2}{\gamma \eta_{pc}} \right]. \quad (22)$$

The total cross section of the hadron interaction which is due to the interaction of slow partons is equal to

$$\begin{aligned} \sigma(\rho_{12}, s_{12}) &= \int d\eta_1 d\eta_2 d^2 \rho'_{12} C(\eta_1) C(\eta_2) \\ &\times \int \frac{d^2 \rho}{(\pi \gamma \eta_{pc})^2} \exp \left[-\frac{(\vec{\rho} - \vec{\rho}_1)^2}{\gamma \eta_{pc}} - \frac{(\vec{\rho} - \vec{\rho}_2)^2}{\gamma \eta_{pc}} \right]. \end{aligned}$$

We have taken into account that $\rho'_1 = \rho + \frac{\rho'_{12}}{2}$, $\rho'_2 = \rho - \frac{\rho'_{12}}{2}$, and that the dependence on ρ'_{12} can be neglected in the exponential factor.

After carrying out the integration over ρ , we obtain

$$\sigma(\rho_{12}, s_{12}) = \frac{e^{-\frac{(\vec{\rho}_1 - \vec{\rho}_2)^2}{2\gamma\eta_{pc}}}}{2\pi\gamma\eta_{pc}} \sigma_0 \quad (23)$$

Inserting (22) into (20), we get

$$\begin{aligned} A_1 &= s_{12} \sigma_0 e^{-\frac{\gamma}{4} q^2 \xi}, \\ \xi &= 2\eta_{pc} = \log \frac{s_{12}}{\mu^2}. \end{aligned} \quad (24)$$

We obtained the scattering amplitude corresponding to the exchange by the Pomeranchuk pole with the slope $\alpha' = \gamma/4$, $\sigma_0 = g^2$ where g is the universal coupling constant of the Pomeron and hadron. The amplitude (24) is usually represented by diagram in Fig. 21

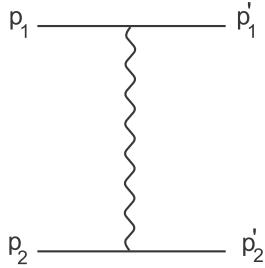


Figure 21:

where a propagator of the form $e^{-\alpha' q^2 \xi}$ corresponds to the Pomeron. In the impact parameter space this propagator has the form (22).

Let us discuss the physical meaning of σ_0 in more detail. For this purpose, let us calculate the zero angle scattering amplitude at ($\vec{q} = 0$), without using the impact parameter representation. The probability of finding a parton with rapidity η and a transverse momentum k_\perp is described by (9). This expression at $\eta \ll \eta_p$ corresponds to the diagram in Fig.22. The wavy line represents integration over parton rapidities from η_p to zero.

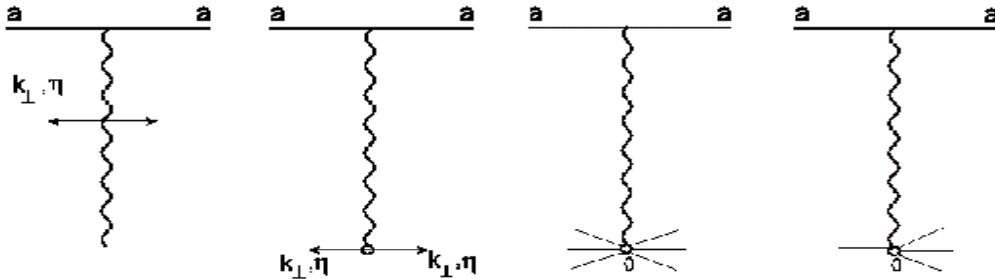


Figure 22

Figure 22a

Figure 22b

Figure 22c

This figure reflects the hypothesis that the calculation of $\varphi(\eta, k_t, \eta_p)$ for sufficiently large η_p and $\eta \ll \eta_p$ leads to an expression for φ which is factorized in the same way as the Pomeron contribution to the scattering amplitude. This is because the parton distribution in this region is independent of the properties of the hadron as well as the values of η, η_p . Compared to the diagram in Fig. 7, Fig. 22 indicates that the calculation of $\varphi(\eta, k_t, \eta_p)$ is similar to the calculation of the inclusive cross section due to the Pomeron exchange. The only difference is that the coupling of the hadron with the Pomeron should be substituted by unity, since a hadron always exists in a Pomeron state. If $\eta \sim 1$, $\varphi(\eta, k_t)$ corresponds to the diagram in Fig.22a, which shows that $\varphi(\eta, k_t)$ depends on η . Similarly, it is possible to determine the probability of finding several slow partons (Fig. 22b), and even the density matrix of slow partons. In this case the amplitude of elastic hadron-hadron scattering in the center of mass frame is determined by

the diagram of fig.23 and the value of σ_0 is determined solely by the interaction of slow partons.

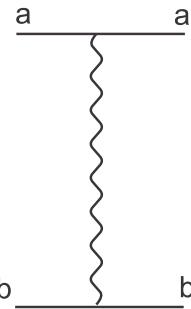


Figure 23

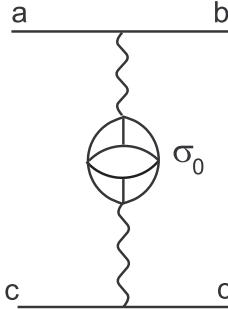


Figure 24

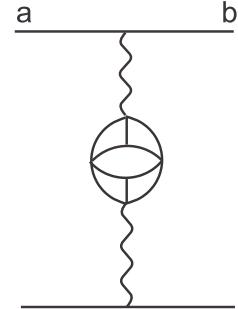


Figure 25

Now let us consider the quasi-elastic scattering, corresponding to the Pomeron exchange (Figs.24,25) at zero transverse momentum. While the probability to find the parton in hadron "a" is determined in eq.(8) by the integral of the wave function squared, the analogous quantity for the amplitude of the inelastic diffractive process (Fig.25) will lead to the integral of the product of the parton functions of different hadrons. They are orthogonal to each other and it is almost obvious that amplitude for inelastic diffractive process at zero angle should vanish for this reason. Indeed the orthogonality condition of eq.(6) has the same structure as the imaginary part of the amplitude. Thus, if at high energies the amplitude factorizes (as it should be for the Pomeron exchange), than the orthonormality condition should also have factorized form in the sense that the integral over parton rapidities with $\eta \ll \eta_p$ factors out, and only constants g_{ab} depend on the properties of specific hadrons (see Fig.26). Orthogonality of the

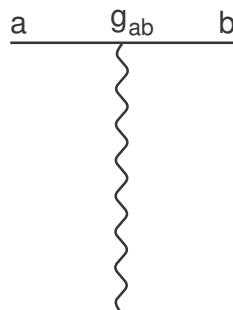


Figure 26:

wave functions of different hadrons implies that $g_{ab} = 0$ at $a \neq b$. In fact the

reason, why the amplitude of inelastic diffractive process vanishes at zero angle is the same as the reason why all cross sections should approach the same value at high energies. Both phenomena are due to the fact that properties of slow partons do not depend on the properties of hadrons to which they belong. We can illustrate this again using the example of quasi-elastic dissociation of the composite system — e.g. deuteron. Let us consider the interaction of a fast nucleon with a deuteron. As we discussed in the previous section, at very large energies partons from different nucleons will always interact with each other independent of the distance between nucleons. This will lead to production of the spectrum of slow partons which does not depend on the relative distance between nucleons in deuteron. This means that the amplitude of the nucleon-nucleon interaction will not depend on the internucleon distance as well. Thus, if nucleons inside the deuteron will remain intact after the interaction, than the deuteron will not dissociate as well, since if the amplitude does not depend on the internucleon distance, the wave function of the deuteron will not change after the interaction.

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